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IRON CORE COLLAPSE MODELS OF TYPE II SUPERNOVAE

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ABSTRACT

Results of recent numerical, one-dimensional core collapse calculations for $10M_{\odot}$, $15M_{\odot}$ and $20M_{\odot}$ population I stars are reviewed. The physics model is discussed, including recent improvements in the nuclear equations of state, and nuclear binding energies. None of the models produces prompt explosions as a direct result of core collapse and bounce.

I. INTRODUCTION

The physics of gravitational collapse of iron cores and its relation to Type II supernovae has been reviewed in this volume by Baym, and by Weaver and Woosley. The physics needed to model core collapse spans more than fifteen orders of magnitude in density ($\rho \lesssim 5\times10^{14}$ g/cm^3), and requires temperatures as high as 30 MeV. For nearly fifteen years, Jim Wilson has played a major role in advancing the frontier of computational work on this problem. During this period, substantial improvements have occurred in our understanding of the conditions under which Type II supernovae are believed to occur (these include improved models of massive stars following silicon burning), and in our understanding of the physics which is important under these new conditions, particularly neutrino interactions with dense matter and the equation of state of nuclear matter.

With each new development, Jim Wilson has modified the one-dimensional (1-D) code to include new physics or improved numerical methods, or to check and calibrate the methods against known results. The collapse code has been discussed by Bowers and Wilson (1982a), who include a complete description of the physics models and numerical algorithms. The evolution of the $10M_\odot$, $15M_\odot$ and $20M_\odot$ iron cores of Weaver, Zimmerman and Woosley (1978) and of Woosley, Weaver and Taam (1980) as described by the 1-D supernova code has been discussed by Bowers and Wilson (1982b). In this article, the physics in the 1-D supernova code will be summarized, including several modifications which were developed subsequent to the discussions of Bowers and Wilson (1982b). Next, the status of prompt explosions from the iron core collapse of stars in the mass range $10M_\odot$ to $20M_\odot$ as described by our most recent calculations will be reviewed. The emphasis here is on prompt explosions, that is, mass ejection immediately following core bounce. In the next article Jim Wilson will discuss the late time behavior of these models.

At the onset of core collapse, the star has an onion-skin structure consisting (from the surface inward) of concentric shells of H, He, nuclei of intermediate atomic weight, Si, and finally iron group nuclei. The latter constitute the iron core. In the recent models of Weaver, Zimmerman and Woosley, and of Woosley, Weaver and Taam, the density in the iron core at the onset of dynamical collapse is a few times 10^9 g/cm^3. The density drops rapidly across the iron core-Si boundary (containing $1.27M_\odot$ to $1.58M_\odot$) to values of order 10^5 g/cm^3. The dynamic time scale at a point in the star at density ρ is proportional to $\rho^{-1/2}$. Thus, although the region outside of the iron core contains most of the stellar mass, the ratio of the dynamic time scale in the core to that in the mantle is of order $(\rho_{mantle}/\rho_{core})^{1/2} \approx 10^{-2}$. Consequently, the core evolution can be considered to be essentially decoupled from the remainder of the stellar model. For our initial model we typically take of the order of $2M_\odot$, which contains the evolved iron core and several tenths $M_\odot$ of the overlying stellar mantle. In all of our calculations, the outer $0.1M_{\odot}$ or so remains nearly stationary well beyond core bounce.

II. IRON CORE COLLAPSE CODE

Our most recent numerical models of core collapse have been constructed using Lagrangian hydrodynamics coupled with the multigroup transport of neutrinos (ν_e, ν_μ, ν_τ and their antiparticles). Analytic and numerically calculated equations of state describe matter for $\rho \lesssim 5\times10^{14}$ g/cm^3 and $T \lesssim 30$ MeV. Thermonuclear burn of carbon, oxygen and silicon are approximated by simple analytic energy release rates. Also included are the emission, absorption, pair production, electron scattering, and coherent nuclear and elastic nucleon scattering of ν_e, ν_μ, and ν_τ and their antiparticles. A simple multigroup model is also used to describe electron capture by an average heavy nucleus. Finally, relativistic corrections in post-Newtonian approximation are included in the momentum equation. The spectral distribution of each neutrino type is calculated dynamically and need not be in thermal equilibrium.

A complete discussion of the numerical algorithms used here can be found in Bowers and Wilson (1982a). The remainder of this section will focus on the physics and the differential equations in the code.

The time rate of change in material energy density results from compression, thermonuclear energy release, and neutrino-matter coupling:

$$\rho\frac{d\varepsilon}{dt} = \frac{P}{\rho}\frac{d\rho}{dt} + \sum_i \rho\varepsilon_i^N X_i + \sum_\alpha \left(\frac{dF^\alpha}{dt}\right)_{\text{collision}} \qquad (1)$$

where ρ is the mass density, ε the specific energy, P the material pressure, ε_i^N are the thermonuclear energy release rates (ergs g^{-1} sec^{-1}) due to C, O and Si burning, X_i are the abundances of C, O and Si, and F^α include the neutrino-matter coupling for all neutrino types (antiparticles included).

The time rate of change of the velocity v is described by

$$\frac{dv}{dt} = -\frac{1}{\rho}\frac{\partial P}{\partial r} + a_G + a_{rot} + a_{rad} \tag{2}$$

where the gravitational acceleration is

$$a_G = -\frac{m(r)G}{r^2}\left[1 + \frac{1}{c^2}\{\varepsilon + P/\rho + \frac{4\pi r^3 P}{m(r)} + \frac{2m(r)G}{r} + \frac{1}{m(r)}\int_0^r[\varepsilon - \frac{m(r')G}{r'}]dm(r')\}\right]. \tag{3}$$

This last expression contains in curly brackets the post-Newtonian correction to the gravitational field of a spherically symmetric mass distribution. Note that m(r) is the baryon mass inside radius r (Zel'dovich and Novikov 1971). The term a_{rot} allows for one-dimensional effects due to rotation, but is not used here. The neutrino radiation acceleration is

$$a_{rad} = \sum a_{rad}^{\alpha} = -\frac{1}{\rho}\sum_{\alpha}\int\left(\frac{D_\nu}{\lambda_\nu c}\right)_\alpha \frac{\partial F_\nu}{\partial r}\,d\nu \tag{4}$$

where the sum is over all neutrino types. For each neutrino type, F_ν^α represents the spectral energy density, and ν is in energy units. The quantity in parentheses reduces to 1/3 in matter which is optically thick to neutrinos, and is included to handle the transition to the optically thin regime properly (see Bowers and Wilson 1982a).

Finally, mass conservation requires that

$$\frac{d\rho}{dt} + \rho\frac{1}{r^2}\frac{\partial}{\partial r}r^2 v = 0\ . \tag{5}$$

The time rate of change of each neutrino distribution includes spacial transport and neutrino-matter coupling.

$$\frac{\partial F^\alpha_\nu}{\partial t} = \frac{1}{r}\frac{\partial}{\partial r} r D^\alpha_\nu \frac{\partial F^\alpha_\nu}{\partial r} + \left(\frac{\partial F^\alpha_\nu}{\partial t}\right)_{\text{collision}} \tag{6}$$

Spacial transport of each neutrino type is described by flux limited spacial diffusion with the diffusion coefficient for neutrinos of type α and energy ν

$$D^\alpha_\nu = \frac{1}{3k^\alpha_\nu + \xi^\alpha_\nu} \quad . \tag{7}$$

Here $k^\alpha_\nu \equiv 1/\lambda^\alpha_\nu$ is the specific opacity, and ξ^α_ν is the flux limiter. The flux limiter is constructed such that $\xi^\alpha_\nu \ll k^\alpha_\nu$ in the diffusion regime. In the optically thin regime $\lambda^\alpha_\nu \gg F^\alpha_\nu \;|\partial r/\partial F^\alpha_\nu|$, ξ^α_ν is chosen such that neutrino energy is transported through matter at the speed of light.

The neutrino mean free paths are given in Bowers and Wilson (1982a). The electron neutrino and antineutrino mean free paths contain contributions for electron-neutrino scattering, neutrino-nucleon scattering, and neutrino coherent scattering off heavy nuclei. The latter process includes finite nuclear structure effects and ion-ion correlation effects. Analytic approximations to the cross sections (Tubbs and Schramm 1975) are used for all processes. The muon and tau neutrino (and antineutrino) mean free paths include contributions from electron scattering, coherent scattering off heavy nuclei, and scattering off free nucleons. Coherent scattering includes ion-ion correlations and finite nucleon structure effects. The cross-sections (Tubbs and Schramm 1975) are used for all processes. The muon and tau neutrino (and antineutrino) mean free paths include contributions from electron scattering, coherent scattering off heavy nuclei, and scattering off free nucleons. Coherent scattering includes ion-ion correlations and finite nucleon structure effects.

In addition to spacial transport, each neutrino spectral energy distribution will change as a result of the coupling between the neutrinos and matter. These are included in the last term on the right hand side of equation (6). For electron neutrinos and antineutrinos these terms correspond to electron capture by heavy nuclei, electron scattering, emission and absorption by free nucleons, compressional work done on the neutrino fields when the matter is optically thick, and finally energy loss associated with radiation acceleration of matter described by equation (4). For muon and tau neutrinos and their antiparticles, the last term on the right hand side of equation (6) includes contributions due to electron scattering, thermal and plasma pair production, compressional heating, and energy changes due to radiation acceleration of matter.

All electron scattering energy exchange processes are described in the Fokker-Planck approximation, which has been calibrated to the Monte Carlo results of Tubbs et al. (Tubbs, Weaver, Bowers, Wilson and Schramm 1980):

$$\left(\frac{\partial F_\nu}{\partial t}\right)_{F.P.} = \nu\frac{\partial}{\partial\nu}\left\{K^\alpha_\nu\left[F^\alpha_\nu(1-f^\alpha_\nu) + kT\left(\frac{\partial F^\alpha_\nu}{\partial\nu} - \frac{3F^\alpha_\nu}{\nu}\right)\right]\right\} \qquad (8)$$

with $f^\alpha_\nu \equiv 2\pi^2\left(\frac{hc}{k}\right)^3 \frac{F^\alpha_\nu}{\nu^3}$,

ν has units of energy, and k is Boltzmann's constant. The analytic approximations to the cross sections used in the diffusion coefficient K^α_ν are discussed in Bowers and Wilson (1982a).

In steady state, equation (8) can be used to show that F^α_ν reduces to the usual Fermi-Dirac distribution $F^\alpha_{\nu,eq}$, whose chemical potential μ_ν is fixed by the constraint $a(\mu_\alpha) = \int F^\alpha_{\nu,eq}\, d\nu/\nu$.

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$$\left(\frac{dF_\nu^e}{dt}\right)_A = n_e c n_p^* \sigma_e(\nu)(1-f_\nu^e) - \frac{cF_\nu^e}{\nu^3} n_n^* \sigma_e(\varepsilon_e)(1-f_e) \ . \qquad (9)$$

Here f_e is the electron distribution function, ε_e the electron energy, n_e the electron number density, and σ_e the electron capture cross-section. Finally, n_p^* and n_n^* represent the effective number density of protons and neutrons, respectively, in the average heavy nucleus which are capable of capturing neutrinos. We have not included capture of $\bar{\nu}_e$.

Similarly, the emission and absorption of electron neutrinos by free nucleons is given by

$$\left(\frac{dF_\nu^e}{dt}\right)_{ea} = n_e c \sigma_{pe} \ n_p \nu(1-f_\nu^e) - c\sigma_{n\nu} \frac{F_\nu^e}{\nu} \ n_n(1-f_e)\nu \quad (10)$$

where n_p and n_n are the number densities of free protons and neutrons. A similar expression is used for electron antineutrinos.

The change in F_ν^α due to radiation acceleration is given by

$$\left(\frac{dF_\nu^\alpha}{dt}\right)_{rad} = \nu\frac{\partial}{\partial\nu} \left[F_\nu^\alpha \frac{1}{r} \frac{\partial}{\partial r} \left(\frac{D_\nu^\alpha}{\lambda_\nu^\alpha c}\right) r^2 v\right] \ . \qquad (11)$$

The quantity in parentheses has been discussed above for equation (6). The change in F_ν^α above has been constructed to conserve the sum of material kinetic energy and neutrino energy per unit volume, to conserve neutrino number for each species, and to conserve total momentum.

and discussed in detail by Bowers and Wilson (1982a).

The neutrino radiation-hydrodynamic equations discussed above require the equations of state

$$P = P(\rho, \varepsilon, X_i, Y_e) \tag{12}$$

and

$$T = T(\rho, \varepsilon, X_i, Y_e) \tag{13}$$

where Y_e is the number of electrons per baryon, and the thermonuclear energy release rates

$$\varepsilon_i^N = \varepsilon_i^N(\rho, T, X_i). \tag{14}$$

The stellar composition includes baryons in the form of free nucleons, and a distribution of atomic nuclei. We assume for simplicity that the baryonic composition can be parameterized by free nucleons, helium nuclei, and an average heavy nucleus of atomic number A and charge Z. The mass fractions of free nucleons X_B, helium, X_{He} and heavy nuclei, X_A, satisfy

$$X_B + X_{He} + X_A = 1. \tag{15}$$

For the purposes of thermonuclear reactions, the average heavy nucleus is represented by a distribution of carbon, oxygen, silicon and :on group nuclei of mass fractions X_C, X_O, X_{Si} and X_{Fe}, respectively:

$$X_A = X_C + X_O + X_{Si} + X_{Fe}.$$

The ratio of free protons to free nucleons is given by

$$Z = \frac{n_p}{n_p + n_n} . \tag{16}$$

Changes in composition due to electron capture, and emission and absorption are described by

$$n_A \frac{dZ_A}{dt} = - \int \left(\frac{dF^e_\nu}{dt}\right)_{ec} \frac{d\nu}{\nu} \tag{17}$$

and

$$\frac{-dn_N}{dt} = \frac{dn_p}{dt} = - \int \left(\frac{dF^e_\nu}{dt}\right)_{ep} \frac{d\nu}{\nu} + \int \left(\frac{d\bar{F}^e_\nu}{dt}\right)_{p\bar{\nu}} \frac{d\nu}{\nu} \tag{18}$$

The subscript ep above denotes the first term on the right hand side of equation (10); subscript $p\bar{\nu}$ denotes the corresponding contribution from the change in $\bar{F}^e_\nu$ due to proton capture of an electron antineutrino.

$\bar{F}^e_\nu$ is the spectral distribution of electron antineutrinos, $n_A = X_A/Am_H$ is the number density of heavy nuclei, $n_e = Y_e\rho/m_H$ and

$$Y_e = (Z/A)X_A + (1/2)X_{He} + Z_B X_B . \tag{19}$$

Finally, the distribution of baryons between heavy nuclei, helium and free baryons is determined by the Saha equation cor-responding to the process

$$(A,Z) \rightarrow \frac{1}{2} Z\ He + (A - 2Z)n$$
$$He \rightarrow 2n + 2p$$

The collapse code is small core contained on a CDC 7600, with an average run time per problem of about 10^2 cycles per minute. Run time to bounce (which is relatively fast) is about five minutes. Following core bounce, the time step drops to several times 10^{-6} sec, and only a few tens of msec can be followed with the fully explicit calculations.

III. RECENT PHYSICS MODIFICATIONS

The equation of state discussed in Bowers and Wilson (1982a, b) has recently been modified, and new collapse models have been calculated. The new models and their effect on core collapse are summarized in this section.

a) Nuclear Equation of State

The numerical equation of state for matter near nuclear density which has been developed by Wilson for core collapse calculations (see Bowers and Wilson 1982b) has been modified to reproduce the analytic results of Bethe, Brown and Cooperstein and Wilson (1983). In this model, the density of nuclear matter $\rho_n = 2.4 \times 10^{14}$ g/cm^3 is about 20% higher than that used in Bowers and Wilson (1982b). The pressure increase due to nuclear repulsion above ρ_n has also been modified such that

$$\delta P_{repulsion} = 7 \times 10^4 \rho(\rho - \rho_n) \text{ dynes/cm}^2 . \qquad (20)$$

Figure 1 shows the new equation of state for $Y_e = 0.30$, and an entropy per baryon $s = k$, where k is Boltzmann's constant. The crosses correspond to the analytic model of Bethe, Brown and Cooperstein and Wilson (1983).

b) Energy of Nuclear Dissociation

The energy needed to thermally dissociate nuclei is an important parameter affecting the strength of the shock when it reaches the mantle. We consider the photodissociation of iron group nuclei into He nuclei and free neutrons

$$(A,Z) \rightarrow (Z/2)\text{He} + (A - 2Z)n$$

The value for the dissociation energy per baryon for He, Q_{He}, used previously has been found to be too large. A better fit can be obtained by noting that changes in density lead to variations in Z and A for our average heavy nucleus. A model for Q_{He} which includes these effects is

$$Q_{He} = 1.7 + 95(0.464 - Z_A)^2 \text{ MeV} . \qquad (21)$$

where the last term takes into account the change in reaction energy with nuclear type. This form grows less rapidly with the charge on the average heavy nucleus Z_A than does the earlier (linear) form. This results in less dissociation of He in the lower density regime outside the homologous core. Consequently, the shock will lose less energy as it traverses the infalling material from the stellar mantle.

c) Results Near Core Bounce

The core model from the Weaver and Woosley $15M_\odot$ star was rerun with the modifications in equation of state and Q_{He} discussed above. The new form of the nuclear equation of state leads to an increase in the sonic mass of 0.02 $M_\odot$, and modification of Q_{He} leads to an additional increase of $0.04M_\odot$. Thus, for the $15M_\odot$ star

$$M_{sonic} \simeq 0.60\ M_\odot. \qquad (22)$$

We note that if general relativistic corrections to the hydrodynamics are turned off, the sonic mass increases to 0.78 $M_\odot$ in agreement with the analytic models of Bethe, Brown, Applegate and Lattimer (1979).

IV. CORE COLLAPSE, BOUNCE AND SHOCK PROPAGATION--PROMPT EXPLOSIONS

Core bounce and shock propagation into the stellar mantle involve essentially all of the physics discussed in Section II. Whether or not the shock can deliver enough energy to the mantle to eject it depends on the net (small) difference between the energy available in the shock and the energy losses as the shock propagates outward.

The initial (gravitational potential) energy of the core at the onset of dynamic infall is of order 10^{49} ergs for most evolved stellar models. The binding energy of the homologous core plus the additional mass which falls in through the shock following bounce is of order 10^{53} ergs. Observations of Type II supernovae indicate that the energy release in the stellar core is of order 10^{50} to 10^{51} ergs. Thus, only about 1% of the energy change available from collapse is needed to produce an explosion.

Analytic models using the results of extensive numerical calculations have been used to clarify the physics of core collapse (see, for example, Brown, Bethe and Baym 1982, and Yahil (1983).

magnitude arguments in an attempt to focus on the basic issues of core collapse models.

Consider the energetics of iron core collapse for low entropy systems such as the Weaver-Woosley models. The initial collapse results in a quasistatic unshocked core of radius $R_c \approx 7 \times 10^6$ cm and $M_H \approx 0.6M_\odot$ (these values are nearly constant for initial stellar masses in the range 10-20 $M_\odot$). Additional material, primarily in the form of heavy nuclei, continues to flow onto M_H. The kinetic energy per baryon of this material at $r \approx R_c$, is of order

$$\frac{1}{2}v_{in}^2 \approx \frac{M_H G}{R_c} \approx 10^{19} \text{ erg/g} \tag{23}$$

which gives $v_{in} \approx 5 \times 10^9$ cm/sec. The shock, which forms just above the surface of the homologous core, must traverse about $\Delta M_\odot \approx 0.4M$ of overlying material before it can reach the neutrinosphere. The infall energy of this mass, ΔM, is to order of magnitude,

$$E_{infall} \approx \frac{1}{2} \Delta M v_{in}^2 \approx 1 \times 10^{52} \text{ ergs.} \tag{24}$$

The kinetic energy of material falling through the shock front goes partly into thermal energy (which helps drive the shock) and partly into neutrino radiation. To order of magnitude the infall energy going into the shock is

$$E_{shock} \approx E_{infall} \approx 10^{52} \text{ erg.} \tag{25}$$

Assuming that this energy goes into particle motion of the matter in ΔM, the energy per baryon is of order $E_{infall}/2\Delta M$, the particle velocity is

$$v \approx (2 \times 10^{52} \text{ergs}/0.8 \times 10^{33} \text{g})^{1/2} \approx 3 \times 10^9 \text{ cm/sec,}$$

and the shock moves outward with velocity $\approx$ 0.1c.

given roughly by

$$\epsilon_B = \frac{m_H GM}{r_D}$$

will be heated to temperatures sufficient to thermally dissociate He and heavy nuclei. For $M = M_\odot$, $r_D = 2 \times 10^7$ cm. Thus the infall energy of matter reaching radii of order r_D or less goes into nuclear thermal dissociation rather than into maintaining the shock's outward motion, and E_{shock} is less than (25). To order of magnitude the shock energy dissipated in this way is

$$E_{diss} = \frac{\Delta M}{m_H} \left(\frac{8\text{MeV}}{\text{baryon}}\right) = 6 \times 10^{51} \text{ ergs.} \tag{26}$$

or about 0.8×10^{19} ergs/g. This corresponds to a fractional reduction in the shock velocity of order

$$[(E_{shock} - E_{diss})/E_{shock}]^{1/2} = v = 0.6.$$

Nuclear recombination subsequently could release this energy, but that is expected to occur on much longer time scales than are of interest for prompt explosions.

Once the shock is within a neutrino mean free path λ_ν of the neutrinosphere, the shock heated matter efficiently converts shock energy into neutrino energy at a rate which is strongly temperature sensitive. At this point the neutrino luminosity changes across the shock front by $\Delta L_\nu = 10^{54}$ erg/sec. The time spent by the shock within a distance λ_ν of the neutrinosphere is of order

$$\Delta t_\nu = \lambda_\nu/v_s = (10^7\text{cm})/(3 \times 10^9\text{cm/sec}) = 3 \times 10^{-3} \text{ sec,} \tag{27}$$

$$E_{rad} = \Delta L_\nu \, \Delta t_\nu = 3 \times 10^{51} \text{ ergs} \tag{28}$$

or about 0.4×10^{19} ergs/g.

Neutrino damping of the shock is seen in all of our recent calculations. For example, Figure 2 shows the change in neutrino luminosity L_ν across the shock, and

$$L_H = 2\pi r^2 \rho v^2 v_s, \tag{29}$$

before the shock reaches the neutrinosphere, and when it is a mean free path beyond the neutrinosphere. Here v is the material velocity, and v_s is the shock velocity.

A prompt explosion appears to be possible only if the shock energy outside the neutrinosphere exceeds 10^{50} to 10^{51} ergs. From the discussion above we see that to within the accuracy of our estimates above

$$E_{shock} - E_{diss} - E_{rad} = 10^{52} - 6 \times 10^{51} - 3 \times 10^{51} \text{ ergs} = 0. \tag{30}$$

The message contained in equation (30) is simple: although core collapse leads to the prompt release of about 10^{52} ergs of gravitational potential energy, it appears to be difficult to convert more than 1% of it into mass motion of the stellar mantle as a direct result of core bounce.

A crude estimate of the shock velocity outside the net ...o-sphere can be made by noting that $L_H \approx 10^{54}$ ergs/sec there; using (29) with $v = v_s$,

$$v_s \approx (L_H/2\pi r_\nu^2 \rho_\nu)^{1/3} \approx 1.0 \times 10^9 \text{ cm/sec}$$

for $\rho_\nu \approx 10^{12}$ g/cm^3 and $r_\nu \approx 10^7$ cm. The escape velocity for matter at r_ν is

$$v_{esc} = [(M_H + \Delta M)G/r_\nu]^{1/2} \approx 3.5 \times 10^9 \text{ cm/sec}$$

which exceeds v_s.

In order to optimize the chances of an explosion, it might be considered sufficient to increase E_{shock} and reduce E_{diss}. Decreasing the stiffness of the nuclear equation of state near core bounce will increase the extent to which the core over shoots its quasi-equilibrium radius. This could impart more energy to the shock. A decrease in leptonization rates during infall can result in an increased homologous core mass. This has two important effects: first, a larger core drives a larger shock at bounce; and second, the greater the mass in M_H, the less overlying mass of heavy nuclei there will be for the shock to photodissociate. Finally, in initial models having higher specific entropy outside the homologous core, less shock energy will be needed to dissociate the heavy nuclei there.

Unfortunately, in all of our calculations we see very little change in the final shock energy due to moderate changes in equation of state, leptonization rates and specific entropy. Although these changes do increase the shock strength, the accompanying increase in shock heating produces an increased rate of neutrino pair production near the neutrinosphere. The net result in all cases that we have investigated is to rob the shock of the extra strength it acquired at bounce.

For example, if the core bounce shock can ultimately reach radii greater than a few times 10^7 cm, no further nuclear thermal dissociation will occur. Furthermore, because of the reduced temperature at larger radii, the neutrino radiative losses will decrease. If energy can be supplied to the shock at this stage, an explosion may be possible (see (Wilson, this volume).

V. IRON CORE COLLAPSE MODELS

Core collapse calculations have been completed for the $10M_\odot$, $15M_\odot$ and $20M_\odot$ models which include the nuclear equations of state and binding energy for He discussed in Section III; the procedure follows that of Bowers and Wilson (1982b). The principal difference between the $15M_\odot$ results of Bowers and Wilson (1982b) and the new model were summarized in Section III above.

None of the models shows any indication of a prompt explosion. In all cases, the shock wave, which is reasonably strong before it reaches the neutrinosphere, turns into what appears to be an accretion shock. The original calculations, which were run out to about 0.03 sec after bounce give no indication that the shock has sufficient energy to accelerate the overlying matter to escape velocities. The results are insensitive to zoning, the exact form of the nuclear equation of state, and the modification in Q_{He} for the $15M_\odot$ model (see Bowers and Wilson 1982b for a discussion of these details).

FIGURE CAPTIONS

(to appear below each figure)

Fig. 1. Nuclear equation of state (solid) including repulsive pressure (20) for the isentrope $s = 1.0$ in units of Boltzmann's constant. The pressure is in dynes/cm^2. The dashed line corresponds to $P \approx \rho^{4/3}$, and the (x) denote values from Bethe, Brown, Cooperstein and Wilson (1983). The arrow denotes nuclear matter density.

Fig. 2. Neutrino damping of shock. a) The neutrino luminosity L_ν and the shock luminosity L_H just before the shock reaches the neutrinosphere. b) L_ν and L_H when the shock is about λ_ν outside the neutrinosphere.

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